

Spinning strings in $AdS_3 \times S^3$ with NS–NS flux

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Abstract

The sigma model describing closed strings rotating in $AdS_3 \times S^3$ is known to reduce to the one-dimensional Neumann–Rosochatius integrable system. In this article we show that closed spinning strings in $AdS_3 \times S^3 \times T^4$ in the presence of NS–NS three-form flux can be described by an extension of the Neumann–Rosochatius system. We consider closed strings rotating with one spin in AdS_3 and two different angular momenta in S^3 . For a class of solutions with constant radii we find the dependence of the classical energy on the spin and the angular momenta as an expansion in the square of the 't Hooft coupling of the theory.

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1. Introduction

Integrability has become a promising path towards a deeper understanding of the AdS/CFT correspondence. After the uncovering of an integrable structure underlying four-dimensional Yang–Mills theory with maximal supersymmetry [1–5], integrability has proved to be a common feature of many other examples of the correspondence. A case of special interest where integrability has also been discovered is the duality between string backgrounds with an AdS_3 factor and maximally supersymmetric two-dimensional conformal field theories, the AdS_3/CFT_2 correspondence. First evidences that integrability could be present in these type of backgrounds

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came from the construction of giant magnon solutions [6]. Later on it was shown that the Green–Schwarz action of type IIB strings with R–R three-form flux compactified on $AdS_3 \times S^3 \times M_4$, where M_4 is either T^4 or $S^3 \times S^1$, is an integrable classical theory [7]. This observation has boosted the analysis of the AdS_3/CFT_2 correspondence using integrability inspired methods [8–24] (for a review see Ref. [25]). Integrability has also been found recently to be a symmetry of more general $AdS_3 \times S^3 \times M_4$ backgrounds with a mixture of R–R and NS–NS fluxes [26]. This discovery has preluded a series of studies of the AdS_3/CFT_2 correspondence with mixed fluxes using integrability [27–32].

Many of the advances in the analysis of $AdS_3 \times S^3 \times M_4$ backgrounds using integrability have been influenced by the developments in the study of the AdS_5/CFT_4 correspondence. In the case of closed string solutions rotating in $AdS_5 \times S^5$ a beautiful picture came from the identification of the corresponding lagrangian with the Neumann–Rosochatius integrable system [33]. A natural question from the point of view of the AdS_3/CFT_2 correspondence is what is the extension of this description to backgrounds with non-vanishing fluxes. This is the problem that we will consider in this note for the case of closed strings rotating in $AdS_3 \times S^3 \times M_4$ with NS–NS three-form flux.

The remaining part of the article is organized as follows. In Section 2 we consider an ansatz for a closed string rotating with two different angular momenta in S^3 and NS–NS three-form flux. The presence of flux introduces an additional term in the Neumann–Rosochatius system. We find a class of solutions with constant radii relying on a similar set of solutions in the absence of flux. In Section 3 we extend the analysis to the case where the string rotates both in AdS_3 and in S^3 . We find a class of solutions with constant radii and one spin in AdS_3 and two different angular momenta in S^3 . We conclude in Section 4 with some general remarks and a discussion on related open problems.

2. Rotating strings in S^3

In this note we will be interested in closed spinning string solutions in $AdS_3 \times S^3 \times T^4$ with NS–NS three-form flux. The solutions that we will study will have no dynamics along the torus and thus we will not include these directions in what follows. The background metrics will then be

$$\begin{aligned} ds_{AdS_3}^2 &= -\cosh^2 \rho dt^2 + d\rho^2 + \sinh^2 \rho d\phi^2, \\ ds_{S^3}^2 &= d\theta^2 + \sin^2 \theta d\phi_1^2 + \cos^2 \theta d\phi_2^2, \end{aligned} \quad (2.1)$$

and the NS–NS B-field will be

$$b_{t\phi} = q \sinh^2 \rho, \quad b_{\phi_1 \phi_2} = -q \cos^2 \theta, \quad (2.2)$$

where $0 \leq q \leq 1$. The value $q = 0$ corresponds to the case of pure R–R flux, where the theory can be formulated in terms of a Green–Schwarz coset. The value $q = 1$ is the limit of pure NS–NS flux, and can be described by a supersymmetric WZW model. In the absence of flux the sigma model for closed strings rotating in $AdS_3 \times S^3$ becomes the Neumann–Rosochatius integrable system, which describes an oscillator on a sphere or a hyperboloid with a centrifugal potential. The presence of flux introduces an additional term in the lagrangian of the Neumann–Rosochatius system [33]. In order to exhibit this it will be convenient to use the embedding

coordinates rather than the global coordinates. The embedding coordinates are related to the global AdS_3 and S^3 angles by¹

$$Y_1 + iY_2 = \sinh \rho e^{i\phi}, \quad Y_3 + iY_0 = \cosh \rho e^{it}, \quad (2.3)$$

$$X_1 + iX_2 = \sin \theta e^{i\phi_1}, \quad X_3 + iX_4 = \cos \theta e^{i\phi_2}. \quad (2.4)$$

In this section we restrict the dynamics of the strings to rotation on S^3 , so that we take $Y_3 + iY_0 = e^{it}$, with $t = \kappa\tau$, and $Y_1 = Y_2 = 0$. For the coordinates along S^3 we choose an ansatz with two different angular momenta,

$$X_1 + iX_2 = r_1(\sigma)e^{i\varphi_1(\tau,\sigma)}, \quad X_3 + iX_4 = r_2(\sigma)e^{i\varphi_2(\tau,\sigma)}, \quad (2.5)$$

where the angles will be taken to be

$$\varphi_i(\tau, \sigma) = \omega_i \tau + \alpha_i(\sigma). \quad (2.6)$$

As we are going to consider solutions that lie on a sphere, the functions $r_i(\sigma)$ must satisfy

$$r_1^2 + r_2^2 = 1. \quad (2.7)$$

As we will be interested in closed strings solutions, the above functions must satisfy

$$r_i(\sigma + 2\pi) = r_i(\sigma), \quad \alpha_i(\sigma + 2\pi) = \alpha_i(\sigma) + 2\pi\bar{m}_i, \quad (2.8)$$

with \bar{m}_i some integer numbers acting as winding numbers. When we enter this ansatz in the Polyakov action

$$S = \frac{\sqrt{\lambda}}{4\pi} \int d^2\sigma [\sqrt{-h} h^{ab} G_{MN} \partial_a X^M \partial_b X^N - \epsilon^{ab} B_{MN} \partial_a X^M \partial_b X^N], \quad (2.9)$$

we find the lagrangian

$$L_{S^3} = \frac{\sqrt{\lambda}}{2\pi} \left[\sum_{i=1}^2 \frac{1}{2} [(r'_i)^2 + r_i^2 (\alpha'_i)^2 - r_i^2 \omega_i^2] - \frac{\Lambda}{2} (r_1^2 + r_2^2 - 1) + q r_2^2 (\omega_1 \alpha'_2 - \omega_2 \alpha'_1) \right], \quad (2.10)$$

where the prime stands for derivatives with respect to σ , Λ is a Lagrange multiplier and we have chosen the conformal gauge. The first piece in (2.10) is the Neumann–Rosochatius integrable system [33]. The presence of the non-vanishing flux introduces the last term in the lagrangian.²

We will now write the equations of motion. The lagrangian is cyclic on the variables α_i . Therefore we easily conclude that

$$\alpha'_i = \frac{v_i + q r_2^2 \epsilon_{ij} \omega_j}{r_i^2}, \quad i = 1, 2, \quad (2.11)$$

where v_i are some integrals of motion and $\epsilon_{12} = +1$ (we assume summation on j). The variation of the lagrangian with respect to the radial coordinates gives us

¹ We will follow closely conventions and notation in [33].

² Note that in the WZW model limit $q = 1$ the lagrangian simplifies greatly because we can complete squares. We will find further evidence on this simplification below.

$$r_1'' = -r_1\omega_1^2 + r_1\alpha_1'^2 - \Lambda r_1, \quad (2.12)$$

$$r_2'' = -r_2\omega_2^2 + r_2\alpha_2'^2 - \Lambda r_2 + 2qr_2(\omega_1\alpha_2' - \omega_2\alpha_1'). \quad (2.13)$$

To these equations we have to add the Virasoro constraints,

$$\sum_{i=1}^2 (r_i'^2 + r_i^2(\alpha_i'^2 + \omega_i^2)) = \kappa^2, \quad (2.14)$$

$$\sum_{i=1}^2 r_i^2 \omega_i \alpha_i' = 0. \quad (2.15)$$

In terms of the integrals v_i the second Virasoro constraint can be rewritten as

$$\omega_1 v_1 + \omega_2 v_2 = 0. \quad (2.16)$$

The energy and the angular momenta of the string are given by

$$E = \sqrt{\lambda} \kappa, \quad (2.17)$$

$$J_1 = \sqrt{\lambda} \int_0^{2\pi} \frac{d\sigma}{2\pi} (r_1^2 \omega_1 - q r_2^2 \alpha_2'), \quad (2.18)$$

$$J_2 = \sqrt{\lambda} \int_0^{2\pi} \frac{d\sigma}{2\pi} (r_2^2 \omega_2 + q r_2^2 \alpha_1'). \quad (2.19)$$

2.1. Constant radii solutions

A simple solution to the equations of motion can be obtained if we take the radii r_i to be some constants, $r_i = a_i$. In this case the derivatives of the angles also become constant and thus

$$\alpha_i = \bar{m}_i \sigma + \alpha_{0i}, \quad (2.20)$$

where the windings become

$$\bar{m}_i \equiv \frac{v_i + q a_2^2 \epsilon_{ij} \omega_j}{a_i^2}. \quad (2.21)$$

The integration constants α_{0i} can be set to zero through a rotation, and the constants \bar{m}_i must be integers in order to satisfy the closed string periodicity condition. The equations of motion for r_i reduce now to

$$\omega_1^2 - \bar{m}_1^2 + \Lambda = 0, \quad (2.22)$$

$$\omega_2^2 - \bar{m}_2^2 - 2q(\omega_1 \bar{m}_2 - \omega_2 \bar{m}_1) + \Lambda = 0, \quad (2.23)$$

and thus we conclude that the Lagrange multiplier Λ is constant on this solution. The Virasoro constraints can then be written as

$$\sum_{i=1}^2 a_i^2 (\bar{m}_i^2 + \omega_i^2) = \kappa^2, \quad (2.24)$$

$$\bar{m}_1 J_1 + \bar{m}_2 J_2 = 0. \quad (2.25)$$

We now find the energy as a function of the angular momenta and the integer numbers \bar{m}_i . In order to do this we first use Eqs. (2.7) and (2.15) to write the radii as functions of ω_i and \bar{m}_i ,

$$a_1^2 = \frac{\omega_2 \bar{m}_2}{\omega_2 \bar{m}_2 - \omega_1 \bar{m}_1}, \quad a_2^2 = \frac{\omega_1 \bar{m}_1}{\omega_1 \bar{m}_1 - \omega_2 \bar{m}_2}. \quad (2.26)$$

With these relations at hand and the definitions (2.17)–(2.19), together with (2.24), we find

$$E^2 = \frac{(J_1 + \sqrt{\lambda} q a_2^2 \bar{m}_2)^2}{a_1^2} + \frac{(J_2 - \sqrt{\lambda} q a_2^2 \bar{m}_1)^2}{a_2^2} + \lambda(a_1^2 \bar{m}_1^2 + a_2^2 \bar{m}_2^2), \quad (2.27)$$

or after some immediate algebra,

$$E^2 = (J_1 + J_2)^2 + J_1 J_2 \frac{(1-w)^2}{w} - 2\sqrt{\lambda} q \bar{m}_1 (J_1 w + J_2) + \lambda(\bar{m}_1 \bar{m}_2 - q^2 \bar{m}_1^2 w) \frac{\bar{m}_1 - \bar{m}_2 w}{\bar{m}_2 - \bar{m}_1 w}, \quad (2.28)$$

where we have made use of (2.25) and we have introduced $w \equiv \omega_1/\omega_2$. Now we need to write the ratio w as a function of the windings \bar{m}_i and the angular momenta J_i . This can be done by adding Eqs. (2.18) and (2.19), subtracting Eq. (2.23) from (2.22), and solving the resulting system of equations,

$$[\bar{m}_1 J - \sqrt{\lambda} q \bar{m}_1 (\bar{m}_1 - \bar{m}_2)]w - \bar{m}_2 J - \sqrt{\lambda} (\bar{m}_1 - \bar{m}_2) \omega_1 = 0, \quad (2.29)$$

$$\omega_1^2 - \bar{m}_1^2 - \frac{\omega_1^2}{w^2} + \bar{m}_2^2 + 2q \bar{m}_2 \omega_1 - 2q \bar{m}_1 \frac{\omega_1}{w} = 0, \quad (2.30)$$

where $J \equiv J_1 + J_2$ is the total angular momentum. When we eliminate ω_1 in these expressions we are left with a quartic equation

$$(\bar{m}_1 w - \bar{m}_2)^2 \left[1 - \left(1 - \frac{\sqrt{\lambda}}{J} q (\bar{m}_1 - \bar{m}_2) \right)^2 w^2 \right] + \frac{\lambda}{J^2} w^2 (\bar{m}_1 + \bar{m}_2) (\bar{m}_1 - \bar{m}_2)^3 (1 - q^2) = 0. \quad (2.31)$$

Rather than trying to solve this equation explicitly, we can write the solution as a power series expansion in large $J/\sqrt{\lambda}$.³ Out of the four different solutions to (2.31), the only one with a well-defined expansion is

$$w = 1 + \frac{\sqrt{\lambda}}{J} q (\bar{m}_1 - \bar{m}_2) + \frac{\lambda}{2J^2} (\bar{m}_1 - \bar{m}_2) (\bar{m}_1 + \bar{m}_2 + q^2 (\bar{m}_1 - 3\bar{m}_2)) + \dots \quad (2.32)$$

which implies that

³ Alternatively we can solve Eq. (2.31) around the WZW limiting point $q = 1$ to get

$$\omega_1 = \frac{J}{\sqrt{\lambda}} + \frac{\sqrt{\lambda} \bar{m}_1 (\bar{m}_1 + \bar{m}_2) J}{(J + \sqrt{\lambda} \bar{m}_2)^2} (1 - q) + \dots, \\ \omega_2 = \frac{J}{\sqrt{\lambda}} - (\bar{m}_1 - \bar{m}_2) + \frac{(J^2 - \lambda \bar{m}_1 \bar{m}_2) (\bar{m}_1 - \bar{m}_2) + \sqrt{\lambda} \bar{m}_2 J (3\bar{m}_1 - \bar{m}_2)}{(J + \sqrt{\lambda} \bar{m}_2)^2} (1 - q) + \dots,$$

which reduce to (2.33) and (2.34) in the limit where $J/\sqrt{\lambda}$ is large.

$$\omega_1 = \frac{J}{\sqrt{\lambda}} + \frac{\sqrt{\lambda}}{2J} \bar{m}_1 (\bar{m}_1 + \bar{m}_2) (1 - q^2) \left[1 - \frac{\sqrt{\lambda}}{J} q \bar{m}_2 + \dots \right], \quad (2.33)$$

$$\begin{aligned} \omega_2 = & \frac{J}{\sqrt{\lambda}} - q(\bar{m}_1 - \bar{m}_2) + \frac{\sqrt{\lambda}}{2J} \bar{m}_2 (\bar{m}_1 + \bar{m}_2) (1 - q^2) \\ & \times \left[1 - \frac{\sqrt{\lambda}}{J} q(\bar{m}_1 + \bar{m}_2) + \dots \right]. \end{aligned} \quad (2.34)$$

Note that the $\mathcal{O}(\sqrt{\lambda}/J)$ terms and the subsequent corrections in (2.33) and (2.34) are dressed with a common factor of $\bar{m}_1 + \bar{m}_2$ that vanishes for equal angular momenta. We can easily prove the existence of this factor if we set $\bar{m}_1 = -\bar{m}_2 \equiv m$ in Eq. (2.31), which reduces to

$$(1 + w)^2 [(J - 2\sqrt{\lambda}qm)^2 w^2 - J^2] = 0, \quad (2.35)$$

whose only well-defined solution is

$$w = \frac{J}{J - 2\sqrt{\lambda}qm}, \quad (2.36)$$

and therefore we can calculate the frequencies ω_1 and ω_2 exactly,

$$\omega_1 = \frac{J}{\sqrt{\lambda}}, \quad \omega_2 = \frac{J}{\sqrt{\lambda}} - 2qm. \quad (2.37)$$

An identical reasoning can be employed to prove the existence of the global factor $1 - q^2$.

If we substitute the value of w in Eq. (2.32) in relation (2.28) we find

$$E^2 = J^2 - 2\sqrt{\lambda}q\bar{m}_1 J + \frac{\lambda}{J} [(\bar{m}_1^2 J_1 + \bar{m}_2^2 J_2)(1 - q^2) + q^2 \bar{m}_1^2 J] + \dots \quad (2.38)$$

When the flux vanishes this expression becomes the expansion for the energy in the Neumann–Rosochatius system describing closed string solutions rotating with two different angular momenta [33]. We must note that the subleading terms not included in (2.38) contain a common factor of $\bar{m}_1 + \bar{m}_2$. Therefore if we look at the particular case of $\bar{m}_1 = -\bar{m}_2$ relation (2.38) simplifies to

$$E^2 = J^2 - 2\sqrt{\lambda}qmJ + \lambda m^2. \quad (2.39)$$

This is the expression for the energy in the case of circular string solutions with two equal angular momenta found in [28]. We stress that relation (2.39) is an exact result because the ratio w is given by Eq. (2.36).

In a similar way if we focus on the case of pure NS–NS flux, where $q = 1$, the energy can also be found exactly,

$$\begin{aligned} E^2 = & \frac{1}{J} \left[J^3 + (\sqrt{\lambda}\bar{m}_1 - J_2)J_1^2 - (\sqrt{\lambda}\bar{m}_1 + J_1)J_2^2 - 2\sqrt{\lambda}\bar{m}_1 J^2 + \lambda\bar{m}_1^2 J \right. \\ & \left. - \sqrt{\lambda}(\bar{m}_1 - \bar{m}_2)J_1 J_2 - \frac{(\sqrt{\lambda}\bar{m}_1 + J_1)(\sqrt{\lambda}\bar{m}_1 - J_2)J^2}{J - \sqrt{\lambda}(\bar{m}_1 - \bar{m}_2)} \right]. \end{aligned} \quad (2.40)$$

3. Rotating strings in $AdS_3 \times S^3$

We now extend the analysis in the previous section to the case where the string can rotate both in AdS_3 and S^3 , again with no dynamics along T^4 . The string solutions that we consider will therefore have one spin S in AdS_3 and two angular momenta J_1 and J_2 in S^3 . We can describe these configurations with the ansatz (2.5)–(2.6), together with

$$Y_3 + iY_0 = z_0(\sigma)e^{i\phi_0(\tau,\sigma)}, \quad Y_1 + iY_2 = z_1(\sigma)e^{i\phi_1(\tau,\sigma)}, \quad (3.1)$$

where the angles are

$$\phi_a(\tau, \sigma) = w_a\tau + \beta_a(\sigma), \quad (3.2)$$

together with the periodicity conditions

$$z_a(\sigma + 2\pi) = z_a(\sigma), \quad \beta_a(\sigma + 2\pi) = \beta_a(\sigma) + 2\pi\bar{k}_a, \quad (3.3)$$

with $a = 0, 1$. Note however that the time direction has to be single-valued so we need to exclude windings along the time coordinate. Therefore we must take $\bar{k}_0 = 0$. When we substitute this ansatz in the Polyakov action in the conformal gauge we obtain again lagrangian (2.10) for rotation in the S^3 piece, together with the contribution from AdS_3 ,

$$L_{AdS_3} = \frac{\sqrt{\lambda}}{4\pi} \left[g^{ab} (z'_a z'_b + z_a z_a \beta'^2_b - z_a z_a w_b^2) - \frac{\tilde{\Lambda}}{2} (g^{ab} z_a z_b + 1) - 2qz_1^2 (w_0 \beta'_1 - w_1 \beta'_0) \right], \quad (3.4)$$

where we have chosen $g = \text{diag}(-1, 1)$ and $\tilde{\Lambda}$ is a Lagrange multiplier. We now write the equations of motion. As the pieces of the lagrangian describing motion along AdS_3 and S^3 are decoupled the equations of motion for r_i and α_i are given directly by expressions (2.11)–(2.13). In a similar way, the equations of motion for z_a are

$$z''_0 = z_0 \beta'^2_0 - z_0 w_0^2 - \tilde{\Lambda} z_0, \quad (3.5)$$

$$z''_1 = z_1 \beta'^2_1 - z_1 w_1^2 - \tilde{\Lambda} z_1 - 2qz_1 (w_0 \beta'_1 - w_1 \beta'_0), \quad (3.6)$$

and the equations for the angles are

$$\beta'_a = \frac{u_a + qz_1^2 \epsilon_{ab} w_b}{g^{aa} z_a^2}, \quad (3.7)$$

where u_a are some integration constants. To these equations we need to add the constraint

$$-z_0^2 + z_1^2 = -1, \quad (3.8)$$

together with the Virasoro constraints, which are responsible for the coupling between the AdS_3 and the S^3 systems,

$$z_0'^2 + z_0^2 (\beta_0'^2 + \kappa^2) = z_1'^2 + z_1^2 (\beta_1'^2 + w_1^2) + \sum_{i=1}^2 (r_i'^2 + r_i^2 (\alpha_i'^2 + \omega_i^2)), \quad (3.9)$$

$$z_1^2 w_1 \beta'_1 + \sum_{i=1}^2 r_i^2 \omega_i \alpha'_i = z_0^2 \kappa \beta'_0, \quad (3.10)$$

where we have set $w_0 \equiv \kappa$. The spin and the energy in this case are given by

$$E = \sqrt{\lambda} \int_0^{2\pi} \frac{d\sigma}{2\pi} (z_0^2 \kappa + q z_1^2 \beta'_1), \quad (3.11)$$

$$S = \sqrt{\lambda} \int_0^{2\pi} \frac{d\sigma}{2\pi} (z_1^2 w_1 - q z_1^2 \beta'_0), \quad (3.12)$$

and the angular momenta are defined again as in Eqs. (2.18) and (2.19).

3.1. Constant radii solutions

As before a simple solution to these equations can be found when the string radii are taken as constant, $r_i = a_i$ and $z_a = b_a$. In this case the periodicity condition on β_0 and the fact that the time coordinate is single-valued implies

$$\beta'_0 = 0. \quad (3.13)$$

Furthermore the angles can be easily integrated again,

$$\beta'_1 = \bar{k}, \quad \alpha'_i = \bar{m}_i, \quad i = 1, 2, \quad (3.14)$$

and thus the equations of motion reduce to

$$w_1^2 - \bar{k}^2 - \kappa^2 + 2q\kappa\bar{k} = 0, \quad (3.15)$$

$$(\omega_2^2 - \omega_1^2) - (\bar{m}_2^2 - \bar{m}_1^2) - 2q(\omega_1\bar{m}_2 - \omega_2\bar{m}_1) = 0. \quad (3.16)$$

The Virasoro constraints become then

$$b_1^2(w_1^2 + \bar{k}^2) + \sum_{i=1}^2 a_i^2(\omega_i^2 + \bar{m}_i^2) = b_0^2\kappa^2, \quad (3.17)$$

$$\bar{k}S + \bar{m}_1 J_1 + \bar{m}_2 J_2 = 0. \quad (3.18)$$

Using the definitions of the energy and the spin, Eqs. (3.11) and (3.12), together with the constraint (3.8), we can write

$$E = \sqrt{\lambda}\kappa + \frac{S(\kappa + q\bar{k})}{\sqrt{\kappa^2 + \bar{k}^2 - 2q\kappa\bar{k}}}. \quad (3.19)$$

We can use now this expression to write the energy as a function of the spin, the two angular momenta and the winding numbers \bar{k} and \bar{m}_i . As in the previous section we can take the second Virasoro constraint together with the condition that $a_1^2 + a_2^2 = 1$ to find that

$$a_1^2 = \frac{\bar{k}S + \sqrt{\lambda}\omega_2\bar{m}_2}{\sqrt{\lambda}(\omega_2\bar{m}_2 - \omega_1\bar{m}_1)}, \quad a_2^2 = \frac{\bar{k}S + \sqrt{\lambda}\omega_1\bar{m}_1}{\sqrt{\lambda}(\omega_1\bar{m}_1 - \omega_2\bar{m}_2)}. \quad (3.20)$$

Taking these relations into account when adding the angular momenta (2.18) and (2.19) we find a relation between the frequencies ω_1 and ω_2 ,

$$\begin{aligned} & [\bar{k}S + \bar{m}_1 J - \sqrt{\lambda} q \bar{m}_1 (\bar{m}_1 - \bar{m}_2)] \frac{\omega_1}{\omega_2} - (\bar{k}S + \bar{m}_2 J) \\ & - \sqrt{\lambda} (\bar{m}_1 - \bar{m}_2) \omega_1 - \frac{q \bar{k} S (\bar{m}_1 - \bar{m}_2)}{\omega_2} = 0, \end{aligned} \quad (3.21)$$

which extends expression (2.29) to the case of spin in AdS_3 . Combining now Eq. (3.16) with (3.21) we can solve for ω_1 . The result is again a quartic equation,

$$\begin{aligned} & [(\omega_1 + q \bar{m}_2)^2 - (\bar{m}_1^2 - \bar{m}_2^2)(1 - q^2)] [\lambda (\bar{m}_1 - \bar{m}_2) \omega_1^2 + 2\sqrt{\lambda} (\bar{m}_2 J + \bar{k}S) \omega_1 \\ & - ((\bar{m}_1 + \bar{m}_2) J + 2\bar{k}S) J] + (\bar{m}_1 + \bar{m}_2) (\bar{m}_1 J + \bar{k}S)^2 (1 - q^2) = 0. \end{aligned} \quad (3.22)$$

Once we have found the solution to this equation, we can read ω_2 from (3.21) and use then the first Virasoro constrain to calculate κ . But before writing the resulting equation let us first take into account that

$$b_1^2 \omega_1^2 + b_1^2 \bar{k}^2 - b_0^2 \kappa^2 = b_1^2 (2\bar{k}^2 - 2q\kappa\bar{k}) - \kappa^2 = \frac{2\bar{k}S(\bar{k} - q\kappa)}{\sqrt{\lambda(\kappa^2 + \bar{k}^2 - 2q\kappa\bar{k})}} - \kappa^2, \quad (3.23)$$

where we have made use of (3.15). The Virasoro constraint becomes thus a sixth-grade equation for κ ,

$$\frac{4\bar{k}^2 S^2 (\bar{k} - q\kappa)^2}{\lambda(\kappa^2 + \bar{k}^2 - 2q\kappa\bar{k})} = (\kappa^2 - a_1^2 (\omega_1^2 + \bar{m}_1^2) - a_2^2 (\omega_2^2 + \bar{m}_2^2))^2. \quad (3.24)$$

The solution to this equation provides κ , and thus the energy, as a function of the spin, the angular momenta, and the winding numbers \bar{k} and \bar{m}_i . However Eqs. (3.22) and (3.24) are difficult to solve exactly. As in the previous section, instead of trying to find an exact solution we can write the solution in the limit $J_i/\sqrt{\lambda} \sim S/\sqrt{\lambda} \gg 1$. Out of the four different solutions to (3.22), the only one with a well-defined limit is

$$\omega_1 = \frac{J}{\sqrt{\lambda}} + \frac{\sqrt{\lambda}}{2J^2} (\bar{m}_1 + \bar{m}_2) (\bar{m}_1 J + \bar{k}S) (1 - q^2) \left[1 - \frac{\sqrt{\lambda}}{J} q \bar{m}_2 + \dots \right]. \quad (3.25)$$

Using now relation (3.21) we find⁴

$$\begin{aligned} \omega_2 &= \frac{J}{\sqrt{\lambda}} - q (\bar{m}_1 - \bar{m}_2) + \frac{\sqrt{\lambda}}{2J^2} (\bar{m}_1 + \bar{m}_2) (1 - q^2) \\ &\times \left[\bar{m}_2 J + \bar{k}S - \frac{\sqrt{\lambda}}{J} q \bar{m}_2 (\bar{m}_1 J + \bar{m}_2 J + 2\bar{k}S) + \dots \right]. \end{aligned} \quad (3.26)$$

Next we can calculate the radii a_1 and a_2 using (3.20), and solve Eq. (3.24) to get

$$\begin{aligned} \kappa &= \frac{J}{\sqrt{\lambda}} - q \left(\bar{m}_1 + 2 \frac{\bar{k}S}{J} \right) + \frac{\sqrt{\lambda}}{2J^2} (\bar{m}_1^2 J_1 + \bar{m}_2^2 J_2 + 2\bar{k}^2 S) (1 - q^2) \\ &- \frac{\sqrt{\lambda}}{J^3} 2q^2 \bar{k}S (\bar{m}_1 J + \bar{k}S) + \dots \end{aligned} \quad (3.27)$$

This result can now be substituted in relation (3.19) to obtain

⁴ Note that as in the case of rotation just in the sphere the $\mathcal{O}(\sqrt{\lambda}/J)$ terms and the subsequent corrections in the expansions for ω_1 and ω_2 are again proportional to $\bar{m}_1 + \bar{m}_2$. We can prove the existence of this factor as in the previous section by setting $\bar{m}_1 = -\bar{m}_2$ in Eq. (3.22).

$$\begin{aligned}
E = J + S - \sqrt{\lambda} q \bar{m}_1 + \frac{\lambda}{2J^2} (\bar{m}_1^2 J_1 + \bar{m}_2^2 J_2 + \bar{k}^2 S) (1 - q^2) \\
+ \frac{\lambda}{J^3} 2q^2 \bar{k}^2 S (J + S) + \dots
\end{aligned} \tag{3.28}$$

In the absence of flux this expression reduces to the expansion for the energy in the Neumann–Rosochatius system for a closed circular string of constant radius rotating with one spin in AdS_3 and two different angular momenta in S^3 [33].

4. Concluding remarks

In this letter we have studied closed string solutions rotating in $AdS_3 \times S^3 \times T^4$ with NS–NS three-form flux. The corresponding string sigma model reduces to the Neumann–Rosochatius integrable system with an additional contribution coming from the non-vanishing flux term. We have considered the cases where the string can rotate either in S^3 with two different angular momenta, or in $AdS_3 \times S^3$ with one spin and two different angular momenta. The equations of motion can be easily integrated either as a power series in $J/\sqrt{\lambda}$ or as a power series around the pure NS–NS point $q = 1$ for the case of constant radii strings. We have found the classical energy in terms of the conserved quantities and the parameter governing the strength of the NS–NS flux.

There are many natural extensions of our analysis in this note. An immediate one is the choice of an ansatz where the worldsheet coordinates τ and σ are exchanged. This kind of ansatz, where the radial coordinates depend on the time variable, corresponds to the pulsating string solutions considered in [34,35]. A similar solution can indeed be readily constructed also in the presence of non-vanishing three-form flux. An equally straightforward continuation of our analysis is the analysis of more general solutions. It would be for instance very interesting to find elliptic solutions with non-vanishing flux, following the analysis in the case of the Neumann–Rosochatius system [33].

Another important question is the study of the spectrum of small quadratic fluctuations [36, 33] around the circular solutions that we have constructed in this note. This problem is a necessary step in order to determine the conditions of stability of our solutions and to find the spectrum of excited string states. It would also be very interesting to extend the analysis in this note to deformations of the AdS_3 backgrounds. An appealing case is that of the η -deformation of $AdS_5 \times S^5$ [37], which has been recently shown to lead to an integrable extension of the Neumann–Rosochatius system [38]. However the η -deformed Neumann–Rosochatius system is much more involved than the one that we have obtained in this note because the deformation is obtained by breaking the isometries of the metric down to the Cartan algebra. Currently only the metric and the NS–NS flux are known for the η -deformation, but if a complete solution was constructed one could expect to be able to introduce a spinning string ansatz with mixed fluxes, depending on the q and η parameters.

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